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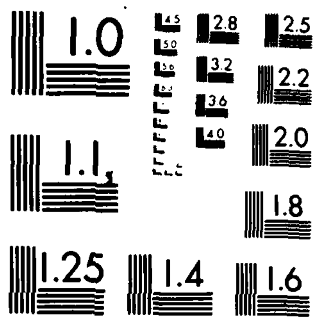

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LECTURES ON MATHEMATICAL COMBUSTION

Lecture 5: Stability of the Plane Deflagration Wave

Technical Report No. 150

J.D. Buckmaster & G.S.S. Ludford

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Theoretical and Applied Mechanics

Thurston Hall
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Cornell University
Ithaca, NY 14853

Contents

Lecture 5

STABILITY OF THE PLANE DEFLAGRATION WAVE

Steady, plane deflagration was introduced in the second lecture; here we shall consider infinitesimal perturbations of it and so examine its stability. We shall find two basic phenomena - the hydrodynamic and Lewis-number effects.

Without the constant-density approximation our task is not easy, because the perturbation equations (though of course linear) have variable coefficients. There are three ways in which this difficulty can be overcome.

(i) If attention is restricted to disturbances whose wavelength is much greater than the thickness of the deflagration, a hydrodynamic description is appropriate (section 3.1). This eliminates the Lewis-number effect, though it may be readmitted as a perturbation (Pelce & Clavin 1982).

(ii) If the constant-density approximation is adopted, then the T- and Y-equations (the only ones that have to be solved) have constant coefficients. This eliminates the hydrodynamic effect, leaving the Lewis-number effect. The hydrodynamics can be reintroduced in the context of a weakly nonlinear theory based on appropriately small heat release.

(iii) For one-dimensional disturbances, the governing equations can be reduced to their constant-density form by means of a von Mises transformation, the distance variable being replaced by the particle function. This requires the diffusion coefficients to be proportional to T (rather than constants), by no means an unreasonable assumption physically.

In this lecture we shall consider the limiting cases (i) and (ii), thereby isolating the hydrodynamic and Lewis-number effects.

1. Darrieus-Landau Instability

The treatment of flame stability from a hydrodynamic viewpoint is due, independently, to Darrieus and Landau. The work of Darrieus, a French aeronautical engineer well known for his invention of the vertical-axis windmill (Darrieus rotor), is often cited as part of the Proceedings of the 1946 International Congress of Applied Mechanics, but these were never published. Copies of a 1938 typescript are in the possession of several members of the combustion community.

As we saw in section 3.1, large-scale disturbances of a plane flame are described by Euler's equations on either side of a temperature discontinuity, namely

$$\nabla\cdot\underset{\sim}{v} = 0,\ \rho D\underset{\sim}{v}/Dt = -\nabla p - \rho g\underset{\sim}{i}. \qquad (1)$$

Bars have been dropped, so it should be remembered that the scales are much bigger than the diffusion time and length. Note that a gravity term has been added, corresponding (when g is positive) to the burnt gas in $x > 0$ lying above the fresh mixture in $x < 0$. The jump conditions (3.9,10) apply across the discontinuity, whose nominal position is $x = 0$.

The undisturbed flow, found as a solution of equations (1) satisfying the jump conditions (3.9,10), is

$$\rho_0 = \begin{cases} 1 \\ 1/\sigma \end{cases},\quad \underset{\sim}{v}_0 = \begin{cases} \underset{\sim}{i} \\ \sigma\underset{\sim}{i} \end{cases},\quad p_0 = \begin{cases} -gx \\ (1-\sigma)-gx/\sigma \end{cases} \quad \text{for } x \lessgtr 0, \qquad (2)$$

where σ is the expansion ratio (3.7). The deformation of the discontinuity is represented by

$$x = F_1(y,t) \qquad (3)$$

if the (small) disturbance parameter is absorbed by F_1 (see figure 1);

we consider perturbations for which

$$F_1 = A \exp(iky + \alpha t) \quad \text{with} \quad k > 0. \tag{4}$$

Restriction to two dimensions, implied by this and

$$\underset{\sim}{v} = (u,v), \tag{5}$$

involves no loss of generality. The goal is to determine the growth parameter α as a function of k, the prescribed wave number of the disturbance. The corresponding perturbations of the flow field are governed by

$$\rho_1 = 0, \quad \underset{\sim}{\nabla}\cdot\underset{\sim}{v}_1 = 0, \quad \rho_0\partial\underset{\sim}{v}_1/\partial t + \rho_0 u_0 \partial\underset{\sim}{v}_1/\partial x = -\underset{\sim}{\nabla}p_1, \tag{6}$$

and the problem is complete once we have found the jump conditions satisfied by these perturbations.

Since only terms that are linear in disturbance quantities are retained, the unit vectors in the directions normal and tangential to the discontinuity are $(1,-F_{1y})$ and $(F_{1y},1)$ so that

$$v_n = u_0 + u_1, \quad v_\perp = u_0 F_{1y} + v_1. \tag{7}$$

The jump conditions (3.9,10) now give

$$u_{1b} = F_{1t}, \; v_{1b}-v_{1f} = (1-\sigma)F_{1y}, \; p_{1b}-p_{1f} = (\sigma^{-1}-1)gF_1 \tag{8}$$

when we note that

$$V = -F_{1t}.$$

Here, following Darrieus and Landau, we have taken

$$W = 1, \text{ i.e. } u_{1f} = F_{1t}, \tag{10}$$

an assumption that can be justified for SVFs with $L = 1$ (section 4.2) and for NEFs (section 4.4). The conditions (8), (10) may all be applied to the undisturbed location of the discontinuity, i.e. $x = 0$.

Solution of the perturbation equations proceeds separately on the two sides of the discontinuity. In the fresh mixture the flow is irrotational (since there are no upstream disturbances), so that

$$u_1 = kP_f e^{kx}, \quad v_1 = ikP_f e^{kx} \quad p_1 = -(\alpha+k)P_f e^{kx} \text{ for } x < 0, \tag{11}$$

if the factor $\exp(iky+\alpha t)$ is omitted. The amplitude P_f of this potential field is as yet unknown. In the burnt mixture there are also rotational terms, due to the generation of vorticity at the curved flame, so that

$$u_1 = kP_b e^{-kx} + kSe^{-\alpha x/\sigma}, \quad v_1 = -ikP_b e^{-kx} - (i\alpha/\sigma)Se^{-\alpha x/\sigma},$$

$$p_1 = (\alpha/\sigma - k)P_b e^{-kx} \quad \text{for} \quad x > 0. \tag{12}$$

In addition to P_b, the amplitude S of the solenoidal field is still to be determined.

The requirements (8),(10) at the discontinuity give four homogeneous equations for A, P_f, P_b, S. These have a non-trivial solution if and only if

$$(\sigma+1)\alpha^2 + 2\sigma k\alpha + (\sigma-1)(g-\sigma k)k = 0, \tag{13}$$

i.e.

$$\alpha = -\frac{\sigma}{\sigma+1} k \pm \sqrt{\frac{\sigma(\sigma^2+\sigma-1)}{(\sigma+1)^2} k^2 - \left(\frac{\sigma-1}{\sigma+1}\right) gk}\ . \tag{14}$$

In the absence of gravity, i.e. for $g = 0$, the larger root is positive for all k, corresponding to instability for all wavelengths. This is the Darrieus-Landau result, which has been an embarrassment to combustion scientists ever since it was discovered. We shall now see how gravity modifies the unacceptable conclusion that plane flames are unstable.

When $g \neq 0$ it is instructive to examine the short and long wavelength limits. As $k \to \infty$,

$$\alpha \sim [-\sigma \pm \sqrt{\sigma(\sigma^2+\sigma-1)}]k/(\sigma+1)+\ldots, \tag{15}$$

so that the influence of gravity dies out and short wavelengths are unstable. As $k \to 0$,

$$\alpha \sim \pm i\sqrt{(\sigma-1)gk/(\sigma+1)} \;-\sigma k/(\sigma+1)+\ldots \tag{16}$$

so that gravity stabilizes long wavelengths. The critical wavenumber separating stable and unstable disturbances is

$$k_c = g/\sigma. \tag{17}$$

Hydrodynamic instability is observed in flames (see, e.g., Sivashinsky 1983), but it is often absent. There are several ways of reconciling this fact with the present results. The flame may be too small to be treated as a hydrodynamic discontinuity; the wavenumbers that allow the disturbed flame to be so treated may be less than k_c; or there may be other stabilizing influences, such as curvature (section 4).

2. The Lewis-Number Effect: SVFs

We now set aside the hydrodynamics and investigate the Lewis-number effect, by adopting the constant-density approximation. The first part of our discussion is concerned with SVFs, which are governed by equation (4.12) with $\Psi = 0$ if volumetric heat losses are negligible.

Section 3.3 found that such flame are unstable to plane disturbances, even finite ones, for $\ell > 1$. Nonplanar disturbances are also unstable, now for $\ell < 1$ as well. This can be demonstrated by examining large-scale disturbances, large even compared to θ. To this end, we introduce new variables

$$(\hat{\eta},\hat{\zeta},\hat{\tau}) = \delta(\eta,\zeta,\tau) \tag{18}$$

and consider the limit $\delta \to 0$. Then $B = O(\delta)$ in the basic equation (4.12), and $W = 1$ to leading order; to next order, we find

$$W = 1 - \tfrac{1}{2}\delta b\hat{K} \quad \text{with} \quad \hat{K} = S^{-1}dS/d\hat{\tau}, \tag{19}$$

an explicit formula for the effect of weak (superficial) stretch on the flame speed. Clearly such stretch decreases the wave speed only for $L > 1$ (cf. section 4.3).

The instabilility of the plane flame for $L < 1$ follows immediately from this relation. Suppose that, because of some disturbance, corrugations have formed (figure 2). The troughs, as viewed from the burnt mixture, experience positive stretch, while the crests are compressed. It follows from the formula (19) with $b < 0$ that the flame speed at the troughs is increased that at the crests decreased, so that the amplitude of the corrugations will grow (corresponding to instability). We are dealing with a cellular instabililty, whose nature is most clearly seen for the NEFs discussed below. Because of the connection (3.27) between the speed and temperature of the flame, the crests are relatively cold and, hence, less luminous than the rest of the flame. For $L > 1$ the effect is reversed, a stable situation.

That this stability conclusion remains valid for disturbances of more moderate scale can also be seen from the basic equation (4.12) of SVFs. Setting

$$W = 1+W_1, \quad K = \kappa_1 \tag{20}$$

gives

$$b(dW_1/d\tau-\kappa_1) = 2W_1 \quad \text{with } W_1 = -F_{1\tau},\ \kappa_1 = F_{1nn} \tag{21}$$

if

$$\chi = -\tau + F_1(\eta, \tau) \tag{22}$$

is the location of the disturbed flame sheet. Growth of the mode proportional to $\exp(\alpha\tau + ik\eta)$ is therefore determined by

$$\alpha^2 - 2b^{-1}\alpha - k^2 = 0, \text{ i.e. } \alpha = b^{-1} \pm \sqrt{b^{-2}+k^2}. \tag{23}$$

For $k \neq 0$, one of the roots is positive whatever the value of b (i.e. L), so that there is instability for all non-planar disturbances when $L \neq 1$. For $k = 0$, there is instability only when b is positive, i.e. for $L > 1$.

The instability of SVFs restricts their value as a frame[illegible]k within which to discuss flames, although the qualitative insights o[illegible]ined from such solutions can be very useful (e.g. concerning flame tip [illegible] see Buckmaster & Ludford 1982 p.159). In pursuit of stable flames, the more usu[illegible]omenon, we now turn to NEFs.

3. The Lewis-Number Effect: NEFs

The equations governing NEFs were derived in section 4.4. To obtain the combusion field of a steady, plane deflagration, we introduce the coordinate

$$n = x + t \tag{24}$$

based on the flame, and seek a solution of the resulting equations

$$\partial T/\partial t + \partial T/\partial n - \partial^2 T/\partial n^2 - \partial^2 T/\partial y^2 = \partial h/\partial t + \partial h/\partial n - \partial^2 h/\partial n^2 - \partial^2 h/\partial y^2 - \ell(\partial^2 T/\partial n^2 + \partial^2 T/\partial y^2) = 0 \tag{25}$$

depending only on n. This yields

$$T_0 = \begin{cases} T_f + Y_f e^n \\ T_b \end{cases}, \quad h_0 = \begin{cases} -\ell Y_f n e^n \\ 0 \end{cases} \quad \text{for } n \lessgtr 0 \tag{26}$$

as the undisturbed temperature and enthalpy (perturbation) profiles,

since equilibrium must prevail behind the flame.

If the equation of the disturbed flame sheet (see figure 1) is

$$n = F_1(y,t), \tag{27}$$

so that the normal derivative in the jump conditions (4.27-29) becomes $\partial/\partial n - F_{1y}\,\partial/\partial y$, then these conditions may be written

$$T_1^- = -Y_f F_1,\ \delta(h_1) = -\ell Y_f F_1,\ \partial T_1^-/\partial n = -Y_f F_1 + h_1^+/\ell_s$$

$$\delta(\partial h_1/\partial n) = \ell h_1^+/\ell_s - 2\ell Y_f F_1 \text{ with } \ell_s = 2T_b^2/Y_f, \tag{28}$$

all quantities being evaluated at the undisturbed flame sheet $n = 0$. (We have used the relation $\phi_* = -h_1^+/T_b^2 + \ldots$ and superscripts $\pm$ to denote values at $n = 0\pm$.) The problem is to solve the equations (25) subject to these conditions and the requirement that T_1, h_1 die out as $n \to \pm\infty$. Arbitrary initial conditions are, as usual, taken into account by considering perturbations for which

$$F_1 = A \exp(iky+\alpha t). \tag{29}$$

The possibility of stability can be seen most easily for $\ell = 0$, when the jump conditions

$$\delta(h_1) = \delta(\partial h_1/\partial n) = 0 \tag{30}$$

ensure that

$$h_1 = 0 \quad \text{for all} \quad n \tag{31}$$

is the appropriate solution of the differential equation (25b). The remaining problem is to solve the T_1-equation for $n < 0$ alone, subject to the boundary conditions

$$T_1 = \partial T_1/\partial n = -Y_1 A \quad \text{at} \quad n = 0; \tag{32}$$

the factor exp(iky+αt) has been omitted. Now, the solutions of the perturbation equation are

$$e^{(1\pm\kappa)n/2} \text{ with } \kappa = \sqrt{1+4\alpha+4k^2} \quad (-\pi/2 < \arg\kappa \leq \pi/2) \tag{33}$$

where, if attention is restricted to unstable modes $Re(\alpha) > 0$, we have

$$Re(\kappa) > 1. \tag{34}$$

It follows that the appropriate solution is

$$T_1 = B\, e^{(1+\kappa)n/2} \quad \text{for} \quad n < 0, \tag{35}$$

and then the boundary condition (32a) requires

$$B = \tfrac{1}{2}(1+\kappa)B, \quad \text{i.e. } \kappa = 1 \quad \text{or } \alpha = -k^2. \tag{36}$$

This contradicts our assumption that the mode is unstable, so we must conclude that there are no unstable modes: when $l = 1$ the flame is stable for all wavenumbers k.

From the results obtained in section 2 for slowly varying disturbances, we may expect that, as ℓ becomes large (positively or negatively), long wavelength disturbances (k small) will become unstable. To see how this happens, we now consider the jump conditions (28) for $\ell \neq 0$. Solution of the perturbation equations proceeds separately on the two sides of the flame sheet. In front we find

$$T_1 = Be^{(1+\kappa)n/2}, \quad h_1 = \{C+\ell\kappa^{-1}[k^2-(1+\kappa)^2/4]Bn\}e^{(1+\kappa)n/2} \quad \text{for } n < 0, \tag{37}$$

where κ has the definition (33b) and again the factor exp(iky+αt) has been omitted. Behind we have

$$T_1 = 0, \quad h_1 = De^{(1-\kappa)n/2} \text{ for } n > 0. \tag{38}$$

The expressions are valid only for $Re(1+\kappa) > 0$ and $Re(1-\kappa) < 0$, i.e. when the condition (34) is satisfied, as it is for unstable modes. The jump conditions (28) now yield a homogeneous system for the coefficients A,B,C,D; a non-trivial solution exists only if

$$2\kappa^2(1-\kappa) + [(1-\kappa)^2-4k^2]\bar{\ell} = 0 \text{ with } \bar{\ell} = \ell/\ell_s, \tag{39}$$

a result due to Sivashinsky.

This dispersion relation should be viewed as determining, for each $\bar{\ell}$, the growth parameter α of any unstable disturbance mode of wave number k. The stability boundaries in the $\bar{\ell}$, k-plane therefore correspond to $Re(\alpha) = 0$; they are shown in figure 3. At any wave number there is a finite band of Lewis numbers, always including $L = 1$, for which the flame is stable and outside of which it is unstable. The left boundary, on which $Im(\alpha)$ vanishes also, is a creation of nonplanar disturbances: the dispersion relation (39) does not yield an α with positive real part for $\bar{\ell} < -1$ and $k = 0$. In fact, the unstable mode becomes neutral as this part of the $\bar{\ell}$-axis is approached. On the other hand, the instability predicted by the right boundary, where $Im(\alpha) \neq 0$, does survive a planar treatment: the dispersion relation (39) with $k = 0$ does yield an α with positive real part for $\bar{\ell} > 2(1+\sqrt{3})$.

The stabilility boundary in the neighborhood of the point $\bar{\ell} = -1$, $k = 0$ can be explained by flame stretch, as was the destabilizing effect of very long wavelength disturbances for $L < 1$ in section 2. We need only note the relation (4.48) between weak stretch and flame speed: for $\bar{\ell} < -1$, positive (negative) stretch decreases (increases) the wave

speed, tending to increase the amplitude of any long-wavelength disturbance (cf. figure 2); for $\overline{\ell} > -1$, amplitudes are decreased.

There are convincing reasons for believing that the left boundary is associated with cellular flames, a common laboratory phenomenon: it corresponds to $\ell < 1$, the Lewis numbers for which such flames are seen; the crests of disturbed flames are darker than the troughs (just as for SVFs with $\ell < 1$), this being a characteristic of cellular flames; and, for $\overline{\ell} < -1$, all modes with $k < \frac{1}{2}\sqrt{-(1+\overline{\ell})}$ are unstable, suggesting that the outcome of the instability will not be monochromatic, another characteristic of cellular flames. In lecture 6 we shall give a nonlinear theory that arises naturally from the present linear analysis and reinforces this conviction.

The right stability boundary is associated with pulsations or traveling waves; it is relatively inaccessible because ℓ is rarely much bigger than 1. However, similar phenomena may be expected in so-called thermites, for which $\ell = \infty$. Otherwise, special means must be devised to make the boundary more accessible. These matters form the subject of lecture 7.

4. The Role of Curvature

The plane premixed flame whose stability has been discussed so far is an idealization seldom approximated, since in practice the flame is usually curved. Even under circumstances designed to nurture a plane state, imperfections can thwart the best efforts and give a curvature, albeit weak, to the flame. In this section we shall investigate certain slightly curved flames, those amenable to the SVF analysis in section 2 and those associated with the left stability boundary for NEFs in section 3.

Consider the cylindrical source/sink flow

$$u = \pm R/r, \quad v = 0, \tag{40}$$

where θr is radial distance and u,v are now polar components. The undisturbed flame is then circular, with θ-multiplied curvature

$$\kappa = \mp R^{-1} \tag{41}$$

determined by the requirement that the mass flux through the unstretched flame is 1. The expressions (20) are replaced by

$$W = 1 + W_1, \quad K = \nabla_{\perp}\cdot v_{1\perp} + \kappa_1 + \kappa W_1, \tag{42}$$

where $v_{1\perp}$ is due to the displacement of the flame in the source/sink flow, which (for the constant-density approximation) is undisturbed. Substitution in the basic equation (4.12) gives

$$b(dW_1/d\tau - \nabla_{\perp}\cdot v_{1\perp} - \kappa_1 - \kappa W_1) = 2W_1 \tag{43}$$

with

$$W_1 = -(R^{-1}F_1 \pm F_{1\tau}), \quad \nabla_{\perp}\cdot v_{1\perp} = \mp 2R^{-2}F_1, \quad \kappa_1 = \pm R^{-2}(F_1 + F_{1\phi\phi}) \tag{44}$$

if

$$r = R + F_1(\phi,\tau) \tag{45}$$

is the location of the disturbed flame sheet (figure 1). The behavior of the mode proportional to $\exp(\alpha\tau + in\phi)$ is therefore given by

$$\alpha^2 - 2(b^{-1} \mp R^{1})\alpha - R^{-1}(\pm 2b^{-1} + n^2R^{-1}) = 0. \tag{46}$$

To discuss this result we introduce the wavenumber

$$k = R^{-1}n \tag{47}$$

and note that the equation (23a) is recovered as $R \to \infty$. There will be instability for those values of k, R that make one of the (always real) roots of equation (46) positive. Such values are bounded by ones for

which $\alpha = 0$, but the reverse is not true. Figure 4 shows the line on which one root vanishes and the corresponding value of the other root; the full part of the line is therefore a stability boundary but the dashed part is not. Results for $R \to \infty$, given in the discussion of equation (23), then show that the stability regions are as lableled. We conclude that curvature is never stabilizing for a sink flow, but that it is for a source flow when

$$\ell < 1, \; k < \sqrt{-2/bR}.$$

We turn now to the neighborhood of $\bar{\ell} = -1, k = 0$ in figure 3, where the dispersion relation (39) reduces to

$$\alpha = -(1+\bar{\ell})k^2 - 4k^4; \tag{48}$$

to have a balanced equation we must require

$$k^2 = O(1+\bar{\ell}), \text{ so that } \alpha = O(1+\bar{\ell})^2. \tag{49}$$

To see how this relation is modified by curvature, we consider once more a circular flame $r = R$ sustained by a source/sink flow. A derivation of the new dispersion relation will be given in section 6.2, where it is needed for a nonlinear analysis. Here we shall only note that the connection between wave speed and stretch, represented by the generalization (49) of the relation (4.48), is preserved; but that recalculations of the stretch, represented by the term in k^2, and of the wave speed, represented by the term α, are needed because of the non-uniformity of the velocity field.

The modified dispersion relation is

$$\alpha = -(1+\bar{\ell})k^2 - 4k^4 \mp R^{-1}, \tag{50}$$

where k is now limited to the values (47); a balancedequation requires

$$R = O(1+\bar{\ell})^{-2} \tag{51}$$

in addition to the earlier restriction (50a). For $\bar{\ell} < -1$ and no curvature, there is a band of wavenumbers

$$0 < k < \tfrac{1}{2}\sqrt{-(1+\bar{\ell})} \tag{52}$$

for which α is positive, corresponding to instability (cf. figure 3). Source flow, i.e. convexity of the flame sheet towards the burnt mixture, narrows the band (figure 5) and, for sufficiently high curvature

$$R^{-1} > (1+\bar{\ell})^2/16, \tag{53}$$

eliminates it. Sink flow, on the other hand, widens the band indefinitely. For $\bar{\ell} > -1$ and no curvature, α is negative for all k, corresponding to stability. Source flow does not change this conclusion, but sink flow produces instability for wavenumbers below an ever-increasing upper limit (figure 5).

For both the SVF and the NEF, convexity towards the burnt mixture is found to be stabilizing, concavity destabilizing.

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Figure Captions

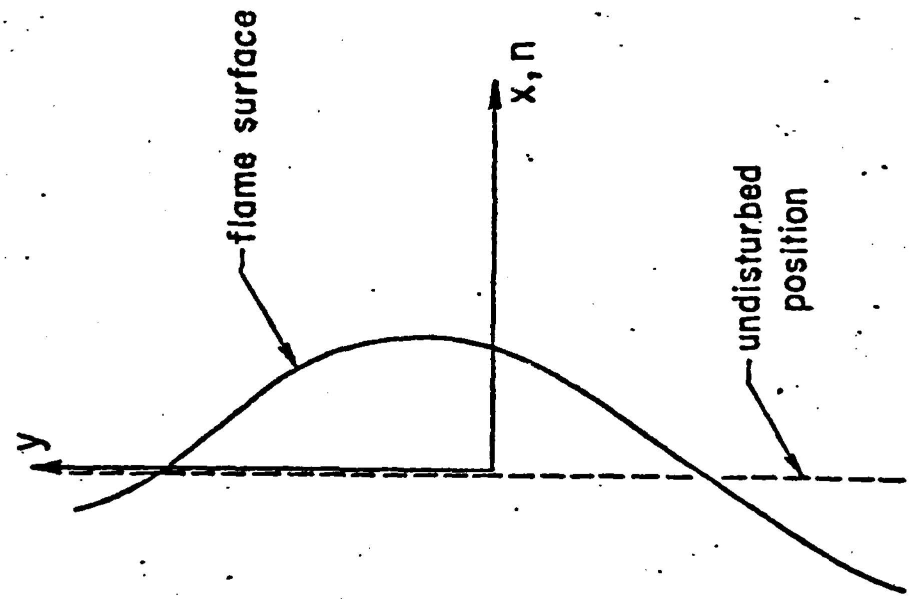
5.1 Notation for stability analysis of plane flame as hydrodynamic discontinuity (x) and as NEF(n). Flame surface is given by $F_1(y,t)$.

5.2 Instability of SVF for $L < 1$.

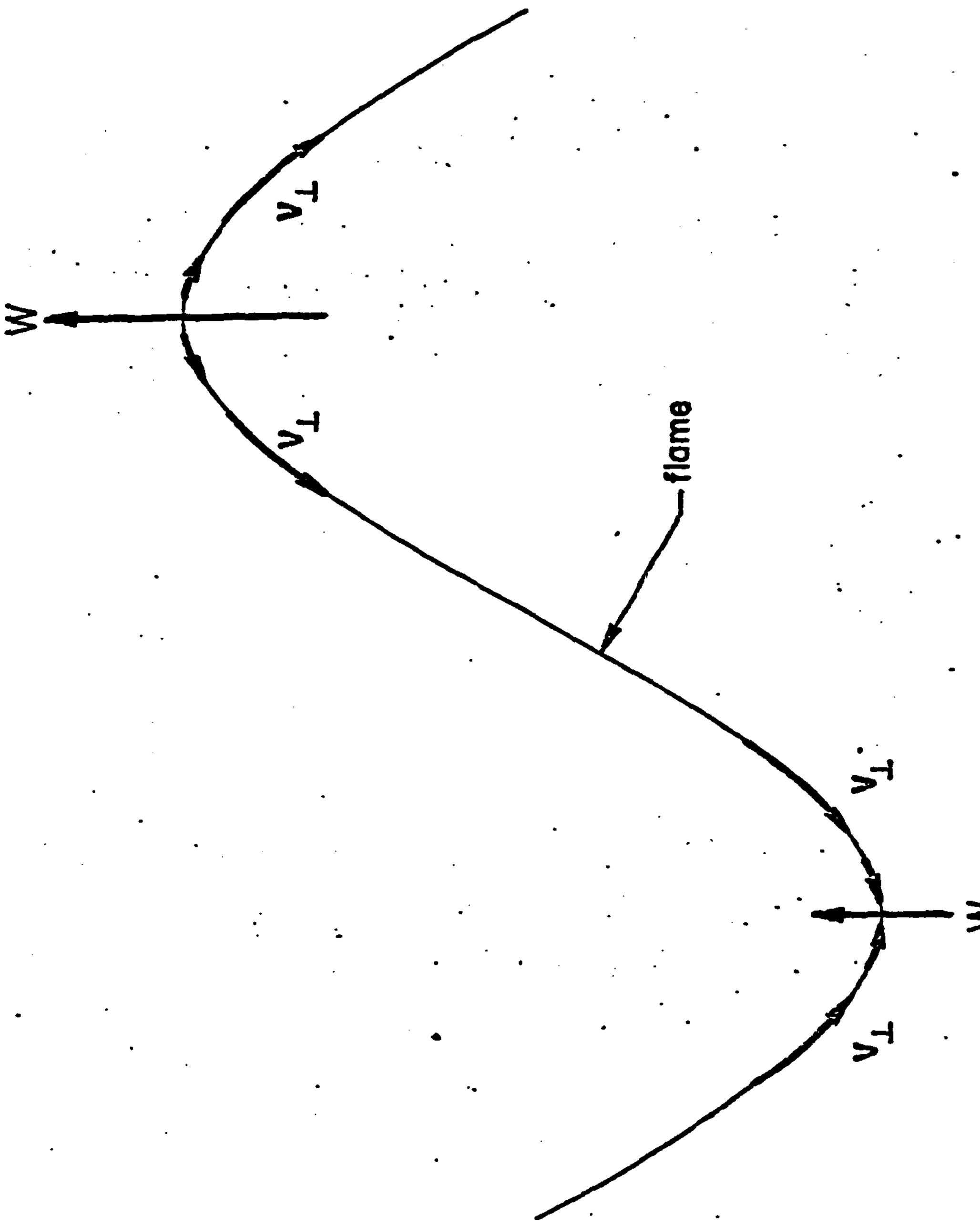
5.3 Linear stability regions for plane NEFs. The boundaries are P: $4k^2 = -(\bar{\ell}+1)$; B: $\bar{\ell} = 2(1+8k^2)$ $[1 + \sqrt{3(1+8k^2)}]/1+12k^2)$.

5.4 Effect of curvature on SVF stability.

5.5 Effect of curvature on NEF stability.


flame surface
x, n
y
undisturbed position

5.2



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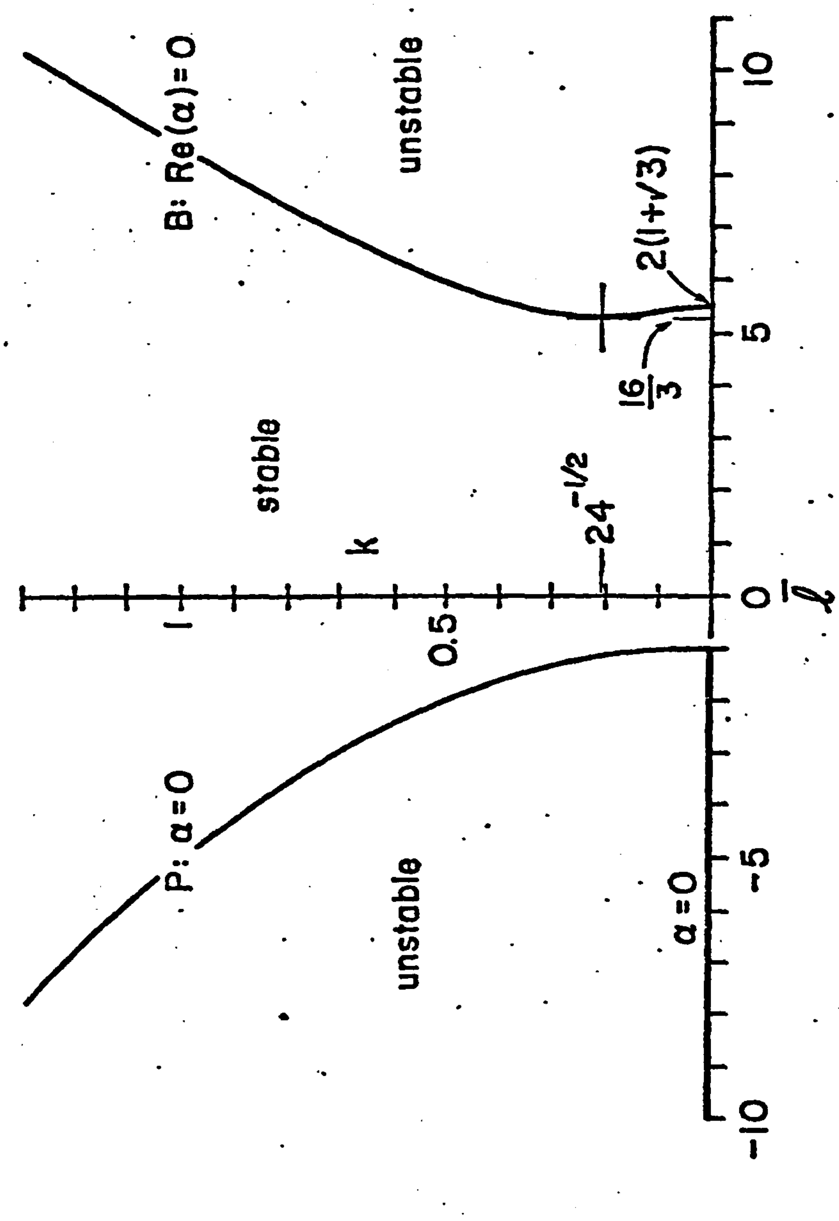
B: Re(α)=0
unstable
stable
P: α=0
unstable
α=0
k
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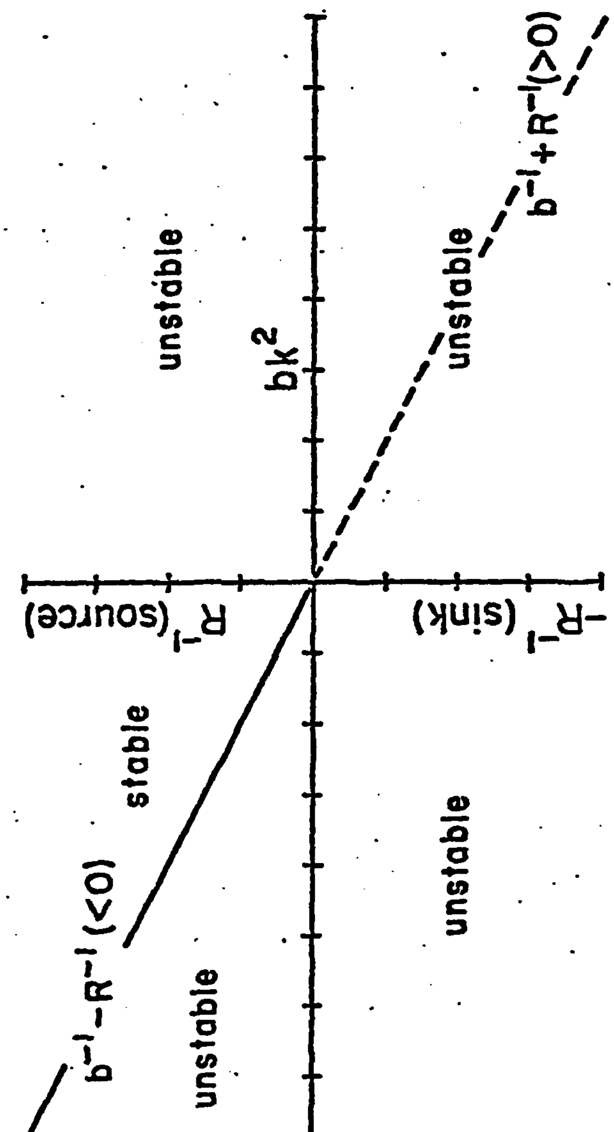
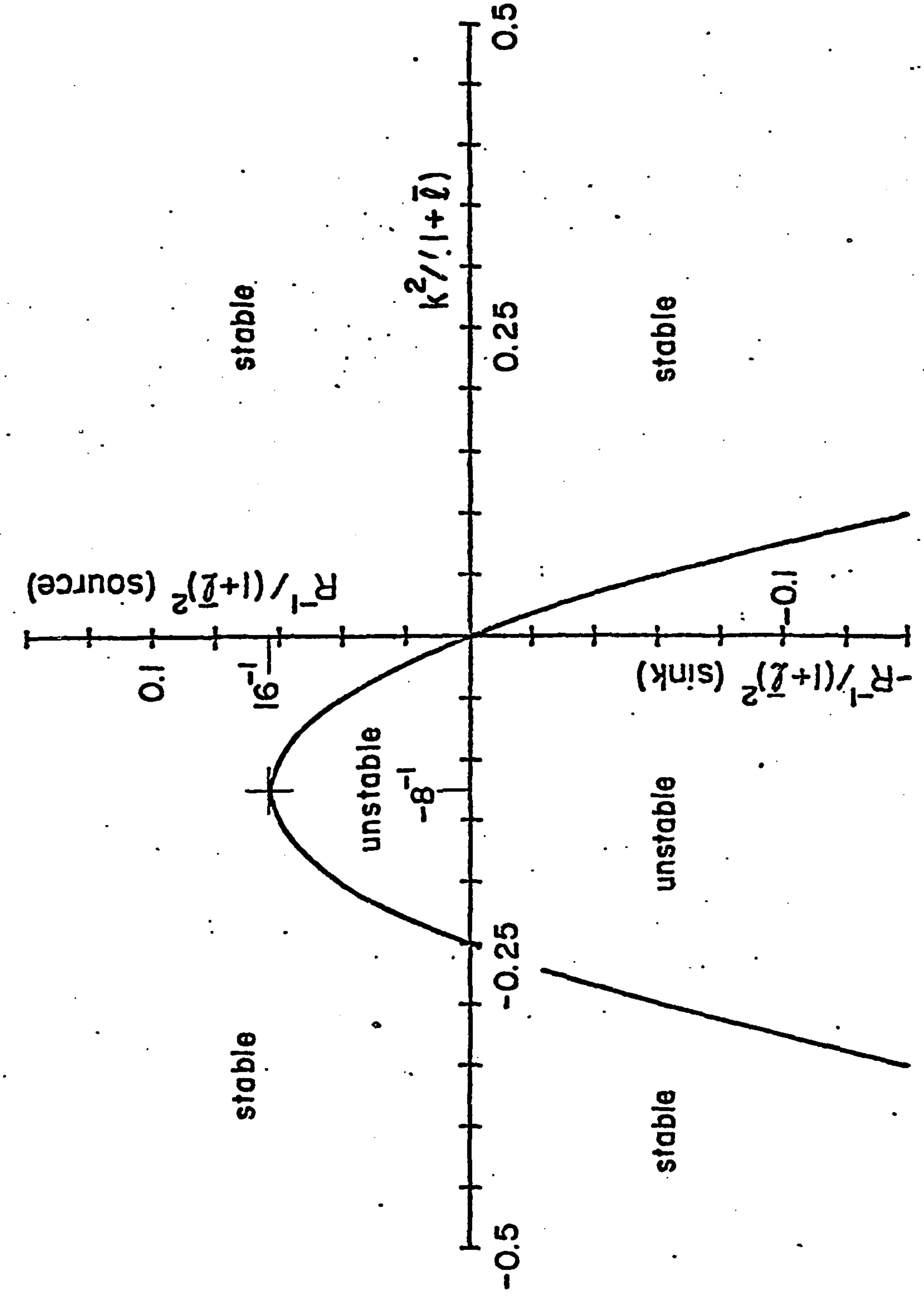
10

5.4



unstable
unstable
$b^{-1}+R^{-1}(>0)$
bk^2
R^{-1}(source)
$-R^{-1}$(sink)
stable
unstable
$b^{-1}-R^{-1}(<0)$
unstable



$k^2/(1+\bar{l})$
$R^{-1}/(1+\bar{l})^2$ (source)
$-R^{-1}/(1+\bar{l})^2$ (sink)
stable
stable
unstable
unstable
stable
stable
0.5
0.25
-0.25
-0.5
0.1
16^{-1}
-8^{-1}
-0.1

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20. ABSTRACT (Continue on reverse side if necessary and identify by block number)

Steady, plane deflagration was introduced in the second lecture; here we shall consider infinitesimal perturbations of it and so examine its stability. We shall find two basic phenomena - the hydrodynamic and Lewis-number effects.

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